Covariant nucleon electromagnetic form factors from the Goldstone-boson-exchange quark model

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Abstract

We present a study of proton and neutron electromagnetic form factors for the recently proposed Goldstone-boson-exchange constituent quark model. Results for charge radii, magnetic moments, and electric as well as magnetic form factors are reported. The calculations are performed in a covariant framework using the point-form approach to relativistic quantum mechanics. All the predictions by the Goldstone-boson-exchange constituent quark model are found in remarkably good agreement with existing experimental data.

PACS: 12.39-x; 13.40.-f; 14.20.Dh

Keywords: Nucleon electromagnetic form factors; Constituent quark model; Point-form quantum mechanics

1 Introduction and motivation

Constituent quark models (CQM) provide a promising tool for the description of low-energy hadron phenomena. They permit the introduction of essential properties of nonperturbative quantum chromodynamics (QCD) and provide a framework for quantitative calculations of hadron properties and reaction observables. In particular, CQM can be made to incorporate the basic symmetries of QCD and, at the same time, of relativistic covariance.

CQM a priori aim at an effective description of QCD rather than its direct solution. They rely on the assumption of constituent quarks (Q), which are viewed as quasiparticles generated dynamically by the spontaneous breaking of chiral symmetry (SBχS). This understanding is in line with the general ideas of dynamical symmetry breaking (e.g., following the ideas of Nambu and Jona-Lasinio [1]) and has recently been further substantiated by lattice QCD results [2].

Over the years there has been an ongoing search for the features of the (effective) Q-Q interaction, specifically in nucleons and more generally in all light as well as strange baryons and their resonances. The debate has been intensified by the appearance of the CQM based on Goldstone-boson-exchange (GBE) dynamics [3]. This type of CQM assumes a linear confinement, as suggested by lattice QCD, with a strength according to the string tension of QCD. For the hyperfine interaction it exploits the idea of GBE [4]. Due to SBχS, the original $\mathrm{SU(3)_L} \times \mathrm{SU(3)_R}$ symmetry of QCD is reduced to an $\mathrm{SU(3)_V}$ (vector) symmetry associated with the appearance of Goldstone bosons. As a result constituent quarks and Goldstone bosons form the essential degrees of freedom governing the low-energy physics of light and strange baryons.

The version of the GBE CQM in Ref. [3] uses a relativistic kinetic-energy operator and advocates, in addition to the linear confinement potential, a pseudoscalar GBE hyperfine interaction; for the latter only the spin-spin component is taken, which is most important in the hyperfine splitting of the baryon spectra. Due to the specific spin-flavor dependence of the hyperfine interaction this kind of CQM has been remarkably successful in reproducing the detailed features of the excitation spectra of all light and strange baryons. Within a single model it has been possible to describe all the resonance levels with the correct orderings. There is no need to separate positive- and negative-parity excitations into different parametrizations, and the particular flavor dependence allows the different level orderings in the N and Λ spectra to be described simultaneously.

Baryon spectroscopy, however, is only a first, though quite demanding, test of low-energy models of QCD. Furthermore, CQM should also provide for a comprehensive description of other hadron phenomena, such as electromagnetic (e.m.) nucleon form factors, resonance excitations and decays, etc. With regard to the GBE CQM the question is whether it can yield as successful a description of these reactions as was the case for spectroscopy. Here we give an answer concerning nucleon e.m. form factors.

In this paper we report first results of proton and neutron electric as well as magnetic form factors calculated with the GBE CQM wave functions in point-form relativistic quantum mechanics. Due to the fact that the GBE CQM uses a relativistic kinetic-energy operator it also provides a covariant mass operator containing interactions given by a Bakamjian–Thomas (BT) construction [5]. Even though the total Q-Q interaction consists of a phenomenological confinement and an instantaneous one-boson-exchange potential, both essentially nonrelativistic, the full Hamiltonian leads to a mass operator fulfilling all necessary commutation relations of the Poincaré group [6]. From the various possibilities for setting up a relativistic quantum theory [7], we have chosen the point-form formulation. It is characterized by several distinctive features. All the dynamics is contained in the four-momentum operators, which commute among themselves and can be simultaneously diagonalized. The generators of the Lorentz boosts contain no interactions and so are purely kinematic; the theory is thus manifestly covariant. In practice, the point-form approach allows us to properly perform Lorentz boosts of the three-Q wave functions and to accurately calculate the matrix elements of the e.m. current operator. In particular, by the introduction of so-called velocity states [8] we can carry out all necessary transformations of the momentum dependences in the wave functions and the relativistic quark spin rotations associated with boosting the nucleon state.

Here we follow the formalism developed in Ref. [9] to investigate the e.m. structure of the nucleons. We calculate the proton and neutron e.m. form factors in point-form relativistic quantum mechanics with one-body currents only. This approach corresponds to a relativistic impulse approximation but specifically in point form. It is called point-form spectator approximation (PFSA) [10] to distinguish it from impulse approximations in other forms of relativistic quantum mechanics, which may lead to different results (see the discussion in Section 3).

In the following we briefly outline the calculations of e.m. current matrix elements with three-Q wave functions in the point-form formulation. The corresponding results for proton as well as neutron electric and magnetic Sachs form factors, charge radii, and magnetic moments are presented in Section 3. A summary and our conclusions are given in Section 4.

2 Nucleon form factors in point form

We start with the eigenstates of the quark-model Hamiltonian for baryons

$$H = \sum_{i=1}^{3} \sqrt{\vec{k}_i^2 + m_i^2} + \sum_{i<j=1}^{3} \left[V^{\mathrm{conf}}(i,j) + V^{\mathrm{hf}}(i,j) \right] . \qquad (1)$$

Herein the first term, with m_i the masses and $\vec{k}_i$ the three-momenta of the constituent quarks, represents the relativistic kinetic energy. The Q-Q interaction consists of the confinement and hyperfine potentials. Specifically we adhere to the version of the GBE CQM as presented in Ref. [3].

The three-Q Hamiltonian (1) is solved using the stochastic variational method (SVM) [11]. This approach yields the eigenenergies with great accuracy; the corresponding results for the ground-state and resonance energies have recently been confirmed by a completely independent method, namely by solving modified Faddeev equations [12]. Moreover, the SVM also produces the three-Q eigenstates in the center-of-momentum frame, i.e. for total nucleon three-momentum $\vec{P} = 0$. The corresponding wave functions can then be interpreted as eigenstates of the mass operator including interactions [9]

$$M = \sqrt{P^\mu P_\mu} = M_{\text{free}} + M_{\text{int}}. \tag{2}$$

Here, $P^\mu = P^\mu_{\text{free}} + P^\mu_{\text{int}}$ is the four-momentum operator with interactions according to the BT construction in point form. We write the general three-Q state defined on the product space $\mathcal{H}_1 \otimes \mathcal{H}_2 \otimes \mathcal{H}_3$ of one-particle spin-$\frac{1}{2}$, positive-mass, positive-energy representations $\mathcal{H}_i = L^2(\mathbb{R}^3) \times S^{1/2}$ of the Poincaré group as

$$|p_1, p_2, p_3; \lambda_1, \lambda_2, \lambda_3\rangle = |p_1, \lambda_1\rangle \otimes |p_2, \lambda_2\rangle \otimes |p_3, \lambda_3\rangle \,, \tag{3}$$

where p_i are the individual quark four-momenta and λ_i the z-projections of their spins. Under general Lorentz transformations U_Λ these states transform as

$$U_\Lambda |p_1, p_2, p_3; \lambda_1, \lambda_2, \lambda_3\rangle = \prod_{i=1}^{3} D^{1/2}_{\lambda'_i \lambda_i}(R_{W_i}) \, |\Lambda p_1, \Lambda p_2, \Lambda p_3; \lambda'_1, \lambda'_2, \lambda'_3\rangle \,. \tag{4}$$

In this equation sums over all $\lambda'_1, \lambda'_2, \lambda'_3$ are understood (here and in the following we adhere to the usual convention of summing over identical indices). The Lorentz transformation in Eq. (4) involves three different Wigner rotations. It is more convenient to first introduce so-called velocity states [8] by applying a particular Lorentz boost $U_{B(v)}$ to the center-of-momentum states, which are defined analogously to Eq. (3) but fulfil the constraint $\vec{P} = \vec{k}_1 + \vec{k}_2 + \vec{k}_3 = 0$,

$$\left| v; \vec{k}_1, \vec{k}_2, \vec{k}_3; \mu_1, \mu_2, \mu_3 \right\rangle = U_{B(v)} \, |k_1, k_2, k_3; \mu_1, \mu_2, \mu_3\rangle \,. \tag{5}$$

Under general Lorentz transformations U_Λ these velocity states transform as

$$\begin{aligned} &U_\Lambda \left| v; \vec{k}_1, \vec{k}_2, \vec{k}_3; \mu_1, \mu_2, \mu_3 \right\rangle = U_\Lambda U_{B(v)} \, |k_1, k_2, k_3; \mu_1, \mu_2, \mu_3\rangle = \\ &= U_{B(\Lambda v)} U_{R_W} \, |k_1, k_2, k_3; \mu_1, \mu_2, \mu_3\rangle = \end{aligned}$$

$$= \prod_{i=1}^{3} D^{1/2}_{\mu_i' \mu_i}[R_W(k_i, R_W)] \left|\Lambda v; R_W \vec{k}_1, R_W \vec{k}_2, R_W \vec{k}_3; \mu_1', \mu_2', \mu_3'\right\rangle , \tag{6}$$

where R_W is the Wigner rotation $R_W(v, \Lambda)$ and $R_W(k_i, R_W)$ is the Wigner rotation of a Wigner rotation. If the boost $B(k_i)$ is chosen to be a canonical one, then $R_W(k_i, R_W) = R_W$ [13], and one has

$$U_\Lambda \left|v; \vec{k}_1, \vec{k}_2, \vec{k}_3; \mu_1, \mu_2, \mu_3\right\rangle =$$
$$= \prod_{i=1}^{3} D^{1/2}_{\mu_i' \mu_i}(R_W) \left|\Lambda v; R_W \vec{k}_1, R_W \vec{k}_2, R_W \vec{k}_3; \mu_1', \mu_2', \mu_3'\right\rangle . \tag{7}$$

Here the Wigner rotations are all the same and the spins can thus be coupled together to a total spin state as in nonrelativistic theory. If a boost other than canonical spin is applied, one can modify Eq. (5) and still arrive at the same transformation property as in Eq. (7). The connection between velocity states and the general three-Q states of Eq. (3) is given by

$$\left|v; \vec{k}_1, \vec{k}_2, \vec{k}_3; \mu_1, \mu_2, \mu_3\right\rangle = \prod_{i=1}^{3} D^{1/2}_{\lambda_i \mu_i}[R_W(k_i, B(v))] \left|p_1, p_2, p_3; \lambda_1, \lambda_2, \lambda_3\right\rangle \tag{8}$$

with $p_i = B(v) k_i$.

In Eqs. (5)–(8) one may equally well express the velocity states in terms of momenta $\vec{p}$ and $\vec{q}$ conjugate to the Jacobi coordinates $\vec{x}$ and $\vec{y}$ (of a certain three-particle partition). The Lorentz transformation in Eq. (7) then specifically reads as

$$U_\Lambda \left|v; \vec{p}, \vec{q}; \mu_1, \mu_2, \mu_3\right\rangle = \prod_{i=1}^{3} D^{1/2}_{\mu_i' \mu_i}(R_W) \left|\Lambda v; R_W \vec{p}, R_W \vec{q}; \mu_1', \mu_2', \mu_3'\right\rangle . \tag{9}$$

Once the Lorentz transformations on the three-Q states have been dealt with, the next task is to calculate the matrix elements of the electromagnetic current operator. For this purpose we follow the formalism developed in Ref. [9]. We first write the invariant nucleon form factors as matrix elements of the one-particle current operator $j^\nu_{[1]}$ (i.e. in the PFSA) in the standard Breit frame, where the momentum transfer Q on the nucleon is only in z-direction $(0, 0, 0, Q) = q_{\rm st}$, as

$$F^\nu_{\mu' \mu}(Q^2) = 3 \int d^3p d^3q d^3p' d^3q' \psi^*_{\mu'}(\vec{p}\,', \vec{q}\,'; \mu_1', \mu_2', \mu_3') \psi_\mu(\vec{p}, \vec{q}; \mu_1, \mu_2, \mu_3)$$
$$\times D^{1/2}_{\lambda_1' \mu_1'}{}^{*}[R_W(k_1', B(v_{\rm out}))] \left\langle p_1', \lambda_1'\right| j^\nu_{[1]} \left|p_1, \lambda_1\right\rangle D^{1/2}_{\lambda_1 \mu_1}[R_W(k_1, B(v_{\rm in}))]$$

$$\times D^{1/2}_{\mu'_2\mu_2}[R_W(k_2, B^{-1}(v_{\rm out})B(v_{\rm in}))] D^{1/2}_{\mu'_3\mu_3}[R_W(k_3, B^{-1}(v_{\rm out})B(v_{\rm in}))]$$
$$\times\delta^3[k'_2 - B^{-1}(v_{\rm out})B(v_{\rm in})k_2]\delta^3[k'_3 - B^{-1}(v_{\rm out})B(v_{\rm in})k_3]. \tag{10}$$

Due to the symmetry of the wave functions it is sufficient to consider only the case where quark 1 is struck by the photon, while quarks 2 and 3 are the spectators, and to multiply the result by 3. The initial and final velocities are given by $m_N v_{\rm in} = p_{\rm st} = (\sqrt{m_N^2 + (Q/2)^2}, 0, 0, -Q/2)$ and $m_N v_{\rm out} = p'_{\rm st} = (\sqrt{m_N^2 + (Q/2)^2}, 0, 0, Q/2)$, respectively, where m_N denotes the nucleon mass.

The single-particle current matrix element in Eq. (10) has the usual form [9]

$$\langle p'_i, \lambda'_i | \, j^\nu_{[1]} \, | p_i, \lambda_i\rangle = e_1 \bar{u}(p'_i, \lambda'_i) \left[\gamma^\nu f_1(\tilde{Q}^2) + i\frac{\sigma^{\nu\rho}\tilde{q}_\rho}{2m_i} f_2(\tilde{Q}^2)\right] u(p_i, \lambda_i) \tag{11}$$

with $u(p_i, \lambda_i)$ the Dirac spinor of quark i and $\tilde{q}_\rho = p'_\rho - p_\rho$, $\tilde{Q}^2 = -\tilde{q}^2$, the momentum transfer on a single quark. It contains the quark invariant form factors f_1 and f_2. In the present study we assume pointlike constituent quarks for which $f_1(\tilde{Q}^2) = 1$ and $f_2(\tilde{Q}^2) = 0$.

From Eq. (10) one obtains the nucleon Sachs form factors through

$$\begin{aligned} F^{\nu=0}_{\mu'\mu}(Q^2) &= G_E(Q^2)\delta_{\mu',\mu} \\ F^{\nu=2}_{\mu'\mu}(Q^2) &= \frac{Q}{m_N} G_M(Q^2)\delta_{\mu',\mu\pm1} \end{aligned} \qquad \mu, \mu' = \pm\frac{1}{2}. \tag{12}$$

Note that only the $\nu = 0$ and $\nu = 2$ components of $F^\nu_{\mu'\mu}$ are needed for the electric and magnetic form factors, respectively. There is no new information in $F^{\nu=1}_{\mu'\mu}$, and one simply recovers G_M. The $\nu = 3$ component of the current vanishes due to current conservation $q^\nu_{st} j_\nu = 0$ (i.e. the continuity equation in the standard Breit frame).

3 Results

The predictions of the GBE CQM [3] for the nucleon e.m. form factors are shown in Fig. 1 below. Their properties at zero momentum transfer are reflected by the charge radii and magnetic moments given in Table 1. The results were calculated in PFSA as explained in the previous section. The input into the calculations consists only of the proton and neutron three-Q wave functions as produced by the solution of the Hamiltonian (1).

One observes that an extremely good description of both the proton and neutron e.m. structure is achieved. It is rather surprising that all relevant observables are quite correctly reproduced. This appears remarkable in view of

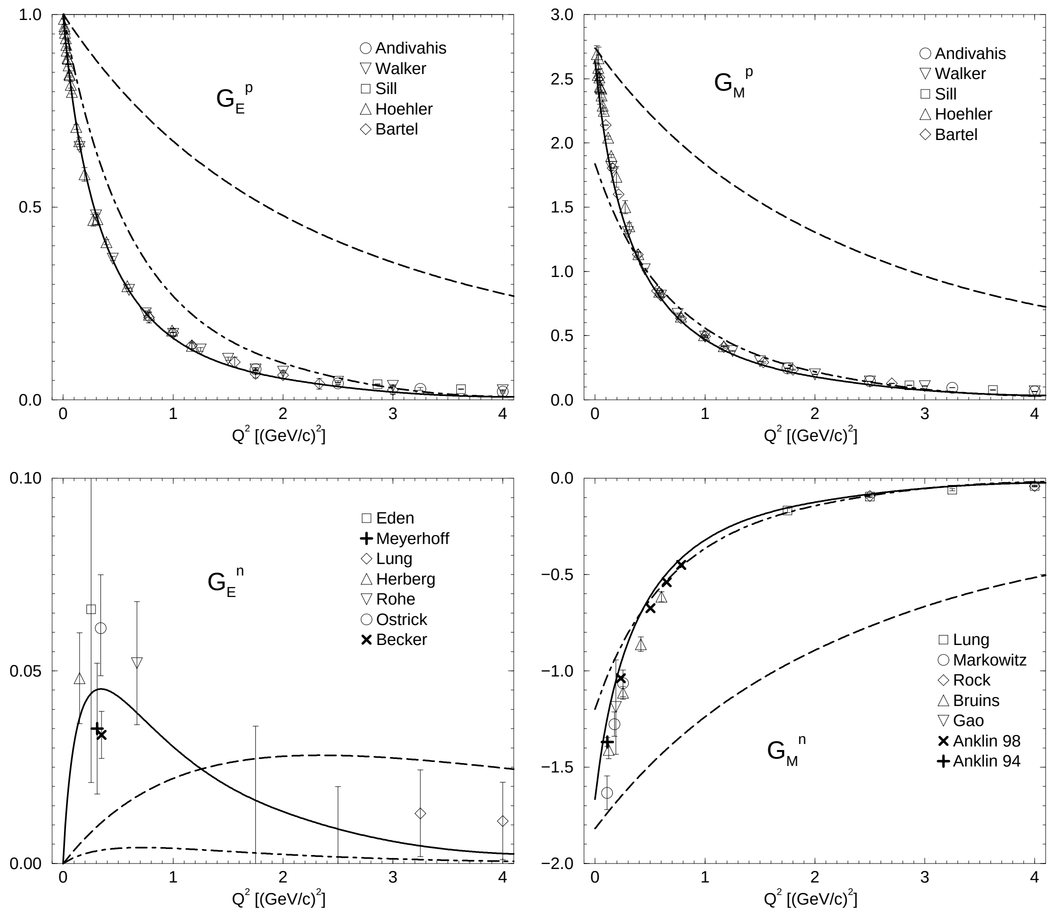

Fig. 1. Proton (upper) and neutron (lower) electric (left) and magnetic (right) form factors as predicted by the GBE CQM [3] in PFSA (solid lines). A comparison is given to the results in NRIA (dashed) and the case with the confinement interaction only (dashed-dotted). The experimental data are from Ref. [14].

the numerous attempts that have been made to explain the low-momentum-transfer e.m. form factors from CQM. Evidently relativity plays a major role here. For comparison we also show results for the form factors when calculated in nonrelativistic impulse approximation (NRIA), i.e. with the standard non-

Table 1
Proton and neutron charge radii as well as magnetic moments as predicted by the GBE CQM [3] in PFSA. A comparison is given also to the results in NRIA and the case with the confinement interaction only.

	PFSA	NRIA	Conf.	Experimental
r_p^2 [fm^2]	0.75	0.10	0.37	0.774(27) [15], 0.780(25) [16]
r_n^2 [fm^2]	−0.12	−0.01	−0.01	−0.113(7) [17]
μ_p [n.m.]	2.64	2.74	1.84	2.792847337(29) [18]
μ_n [n.m.]	−1.67	−1.82	−1.20	−1.91304270(5) [18]

relativistic form of the current operator and no Lorentz boosts applied to the nucleon wave functions. Evidently there is no way of describing the nucleon e.m. form factors in a nonrelativistic theory. However, also in comparison to other relativistic attempts (see, e.g., [19–21]) the covariant results obtained here with the use of realistic CQM wave functions in the point-form approach appear noteworthy, since the e.m. form factors of both the neutron and the proton are readily explained even with pointlike constituent quarks. At least for the range of momentum transfers considered in Fig. 1 there is no need to introduce constituent quark form factors (or any other phenomenological parameters beyond the CQM). This has been necessary in previous relativistic studies in order to bring the theoretical predictions in agreement with experimental data [19–21]. Here, in particular, the momentum dependences already have the right behaviour. E.g., the proton electric form factor nicely matches the dipole form for $Q^2 \lesssim 1$ GeV^2, while it starts to deviate from it beyond, following the trend of recent JLab data [22]. It is only with regard to the magnetic moments that there remains a small difference between the theoretical predictions and the experimental data. An explanation for this gap might be offered by a recent study of pion-loop corrections to magnetic moments [23].

The results presented here depend predominantly on the relativistic boost effects introduced into the nucleon wave functions. The corresponding Lorentz transformations affect the quark spins and the momentum dependences of the wave functions (cf. Eq. (10)). In the point-form approach we are able to perform these transformations without any approximations. The calculations are facilitated by the fact that in point form all interactions are contained in the momentum operators while the generators of the Lorentz group remain interaction-free. This represents an important technical advantage, in that the angular momenta and Lorentz boosts are just the same as in the free case.

In order to get an idea of the role of the GBE hyperfine interaction in the e.m. form factors, we have also considered the case with the confinement potential only. In addition to differences in the wave functions, the nucleons now also have a larger mass of $m_N^{\rm conf} = 1353$ MeV. This different mass is very important for the behavior of the form factors for low Q^2 and is essentially responsible for the corresponding results given in Table 1. Shifting the nucleon mass artificially to $m_N = 939$ MeV would change the charge radii and magnetic moments in the following way: $r^2 \to r^2 \left(m_N^{\rm conf}/m_N\right)^2$ and $\mu \to \mu \left(m_N^{\rm conf}/m_N\right)$. As a result the proton charge radius as well as the magnetic moments of both the proton and the neutron would then already be very close to the values obtained with the full interaction. Only the neutron charge radius would still remain much too small, due to the absence of the mixed-symmetry component in the wave function for the case with the confinement potential only. Though the mixed-symmetry component brought about by the hyperfine interaction is rather small, it turns out to be most essential for reproducing the neutron charge radius in a reasonable manner.

4 Summary and Conclusions

We have presented a first study of elastic nucleon e.m. form factors with the GBE CQM in the point-form approach. The theoretical predictions obtained in PFSA are found to be in remarkably good agreement with all experimental data (charge radii, magnetic moments, electric and magnetic form factors) both for the proton and the neutron. No further ingredients beyond the quark model wave functions (such as constituent quark form factors etc.) have been employed. Only relativistic boost effects are properly included in point-form relativistic quantum mechanics.

Our results suggest that relativistic boost effects are most important in the calculation of nucleon e.m. observables. The point-form approach appears advantageous from a practical point of view, as it makes it possible to include all boost effects in the evaluation of the current-operator matrix elements.

After the successful description of the spectroscopy of all light and strange baryons in a unified framework [3], the GBE CQM now appears capable of also explaining the first dynamical observables, namely the nucleon e.m. structure. It will be important to perform a series of further detailed studies related to the present investigation of proton and neutron elastic form factors. At the same time one is immediately tempted to ask how well electromagnetic transitions (and further on hadronic reactions such as baryon resonance decays) can be described. The point-form approach offers the possibility for performing the relevant investigations on a reliable basis.

Acknowledgements

We have profited from numerous discussions with L.Ya. Glozman. The work was partly performed under the contract ERB FMRX-CT-96-0008 within the frame of the Training and Mobility of Researchers Programme of the European Union.

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